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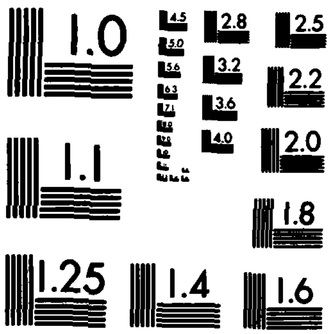
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A Radiative Transfer Model For Undersea Optics Applications

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Measurements of the differential light scattering cross-section in the ocean show that, while the volumetric differential scattering cross section $d\sigma/d\Omega$ is highly peaked in the forward direction, the mean-square scattering angle is not small. Consequently, the many previous theoretical treatments based on the smallness of $<\theta^2>$ are not applicable.

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ABSTRACT

Measurements of the differential light scattering cross-section in the ocean show that, while the volumetric differential scattering cross section $d\beta/d\Omega$ is highly peaked in the forward direction, the mean-square scattering angle

$$< \theta^2 > = (\beta_T)^{-1} \int \theta^2 (d\beta/d\Omega) d\Omega$$

is not small. Consequently, the many previous theoretical treatments based on the smallness of $< \theta^2 >$ are not applicable. The differential scattering cross-section is well approximated by the form

$$\beta(\theta) = \beta_T \theta_0 [\pi^2 \theta (\theta_0^2 + \theta^2)]^{-1}$$

where $\theta_0 \approx .12$ radians (7^o). A solution to the radiative transfer equation is found for this differential cross section. Many useful results can be obtained analytically; e.g., the spatial spreading of an initial pencil beam is given by

$$\Delta r = \theta_0 \int_0^\tau d\tau' [z(\tau) - z(\tau')] ,$$

where Δr is the full width at half maximum, and $z(\tau)$ is the depth associated with optical depth $\tau = \int_0^z \beta_T(z')dz'$. This rate of spreading exceeds that of treatments based on a small mean-square scattering angle.

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INTRODUCTION

The propagation of light in clear ocean waters is dominated by two processes: absorption and scattering by nonabsorbing particles and plankton. Measurements of the differential scattering cross section show that the relative angular dependence of the scattering cross-section is remarkably constant even though the total volumetric scattering cross section β_T varies substantially. This suggests that we represent the volumetric differential scattering cross section $\beta(z,\theta)$ by the separable form

$$\beta(z,\theta) = \beta_T(z)\ X(\theta) \tag{1}$$

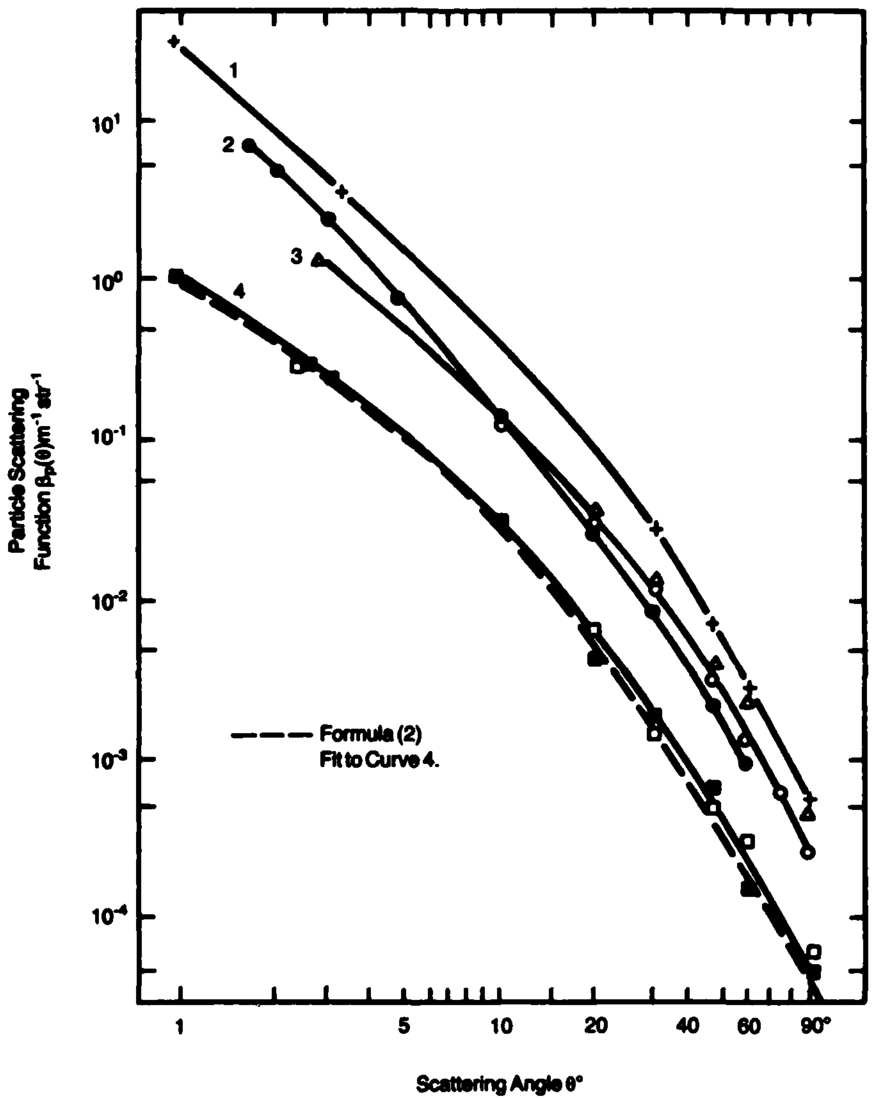
where z denotes depth. Figure 1 shows that the observed angular dependence of the scattering can be well represented by the formula

$$X(\theta) = \frac{\theta_0}{\pi^2 \theta \left(\theta_0^2 + \theta^2\right)} \tag{2}$$

$$\theta_0 = .12 \text{ radians (7 degrees)}$$

where $X(\theta)$ satisfies the small-angle normalization condition

Figure 1. Measurements of volumetric differential scattering cross-section and the theoretical fit, formula (2). [From Fig. 15, N. G. Jerlov, Marine Optics, Elsevier, (1976), p.38.]



$$\int_0^\infty X(\theta) 2\pi\theta d\theta = 1 \; . \tag{3}$$

Formula (2) fits the data well in the region $0 < \theta < \frac{\pi}{2}$; the contribution from angles in the region $\frac{\pi}{2} < \theta < \pi$ to the total scattering cross section is small.

Let us note that, while the normalization integral (3) converges nicely, the mean-square scattering angle

$$< \theta^2 > = \int_0^\infty X(\theta) \; 2\pi\theta^3 d\theta = \int_0^\infty \frac{2\theta_0 \theta^2 d_\theta}{\pi\left(\theta_0^2 + \theta^2\right)} \tag{4}$$

is divergent when small-angle approximations are used. For this reason, the previous[1-3] treatments which were based on a small mean-square scattering angle and a Fokker-Planck diffusion in angles are not appropriate. In this work, we do not use a Fokker-Planck model, but retain the small angle approximation to obtain a direct solution to the radiative transfer equation. A closed expression for the rate of radial spreading of an initial pencil beam is found. For laser-radar applications, the scaling of the range straggling of a pulse can be found analytically, but the actual distribution comes from a numerical solution to a partial differential equation.

RADIATIVE TRANSFER EQUATION

The radiative transfer equation in the small-angle approximation can be written

$$\frac{n}{c}\frac{\partial f}{\partial t} + \frac{\partial f}{\partial z} + \vec{\theta} \cdot \frac{\partial f}{\partial \vec{r}} = -(\alpha+\beta_T)f$$

$$+ \beta_T \int d^2\vec{\theta}' \; f(\vec{\theta}') \; X(\vec{\theta} - \vec{\theta}') \qquad (5)$$

where $\alpha(z)$ is the absorption coefficient, $\beta_T(z)$ is the total scattering cross section, and n is the index of refraction. The variation of α and β with depth z is arbitrary. The solution to (5) takes the form

$$f = g\left(\frac{c}{n}\,t-z\right) \; \psi(z,\vec{r},\vec{\theta}) \qquad (6)$$

where the time dependence of g can be arbitrary and ψ satisfies the time-independent equation.

$$\frac{\partial \psi}{\partial z} + \vec{\theta} \cdot \frac{\partial}{\partial \vec{r}}\,\psi = -(\alpha+\beta_T)\psi + \beta_T \int d^2\vec{\theta}'\,\psi(z,\vec{r},\theta') \; X(\vec{\theta}-\vec{\theta}') \qquad (7)$$

Equation (7) can be solved by Fourier transform. Introducing the definition

$$\psi(z,\vec{r},\vec{\theta}) = \int \frac{d^2\vec{k}}{(2\pi)^2} e^{i\vec{k}\cdot\vec{r}} \int \frac{d^2\vec{p}}{(2\pi)^2} e^{i\vec{p}\cdot\vec{\theta}} \Psi(z,\vec{k},\vec{p}) \tag{8}$$

one arrives at the equation

$$\frac{\partial\Psi}{\partial z} - \vec{k} \cdot \frac{\partial\Psi}{\partial\vec{p}} = -(\alpha+\beta_T)\,\Psi + \beta_T\,\Psi\,X(p)\ . \tag{9}$$

The Fourier transform of the scattering function depends only on the magnitude of p .

$$X(p) = \int d^2\vec{\theta}\ e^{-i\vec{p}\cdot\vec{\theta}}\ X(\theta) = 2\pi \int_0^\infty \theta d\theta \cdot J_0(p\theta)\ X(\theta)$$

$$= \frac{2}{\pi}\int_0^\infty \frac{du}{1+u^2} J_0(p\theta_0 u) = \frac{2}{\pi}\int_0^1 \frac{e^{-p\theta_0 v}}{(1-v^2)^{1/2}}\,dv \tag{10}$$

Sections 6.532, 8.551, and 8.43 of Gradshteyn and Ryzhik[4] were used to transform the integral.

Equation (9) is hyperbolic and is solved by integrating along characteristics

$$\Psi(z,\vec{k},\vec{p}) = \exp\left\{-\tau_\alpha + \int_0^z \beta_T(z')\left[X(s)-1\right]dz'\right\} \tag{11}$$

where

$$\tau_\alpha = \int_0^z \alpha(z')dz \tag{12}$$

$$s = |\vec{p} + \vec{k}\ (z-z')| \tag{13}$$

when $z = 0$, Ψ takes on the value $\Psi = 1$ corresponding to a delta-function beam $\psi = \delta(\vec{r})\ \delta(\vec{\Theta})$. Equations (8), (10), (11)-(13) constitute a solution of the small-angle radiative-transfer problem.

APPLICATIONS

In general, an evaluation of the integrals (8) and (11) calls for numerical work. But fortunately, two important special cases can be evaluated analytically. Our interest centers on the case where the optical depth to small-angle scattering

$$\tau_\beta \equiv \int_0^z \beta_T(z')dz' \tag{14}$$

is large. In this limit, we can utilize the Taylor expansion of X(s)

$$X(s) = 1 - \frac{2}{\pi} s\theta_0 + \frac{1}{4} s^2\theta_0^2 + \ldots . \tag{15}$$

obtainable directly from the last representation of equation (10). Only the first term of this series need be retained, and we find that

$$\Psi(z,\vec{k},\vec{p}) = \exp\left\{-\tau_\alpha - \frac{2}{\pi}\int_0^z \beta_T(z')dz' \,|\vec{p} + \vec{k}\,(z-z')|\right\} \tag{16}$$

The first special case concerns the intensity as a function of radius integrated over all angles

$$I(z,\vec{r}) = \int d\vec{\Theta}\ \psi(z,\vec{r},\vec{\Theta}) = \int \frac{e^{i\vec{k}\cdot\vec{r}}d\vec{k}}{(2\pi)^2}\ \Psi\ (z,\vec{k},\vec{p})\ \Big|_{\vec{p}=0}$$

$$= e^{-\tau_\alpha} \int_0^\infty \frac{kdk\ J_0(kr)}{2\pi} \exp\left(-\frac{2}{\pi} k\Theta_0 \bar{z}\right) \tag{17}$$

where

$$\bar{z} = \int_0^z \beta_T(z')(z-z')dz' = \int_0^{\tau_\beta} d\tau_\beta'\ [z-z(\tau_\beta')]\ . \tag{18}$$

In the special case where β_T is uniform, one finds

$$\bar{z} = \frac{1}{2}\beta_T\ z^2 = \frac{1}{2}\tau_\beta z \tag{19}$$

The integral in (17) can be carried out [Gradshteyn and Ryzhik[4], section 6.623] to give

$$I(z,\vec{r}) = \frac{\bar{z}\Theta_0 e^{-\tau_\alpha}}{\pi^2\left[r^2 + \left(\frac{2}{\pi}\bar{z}\Theta_0\right)^2\right]^{3/2}} \tag{20}$$

Equation (20) governs the spreading of an initial pencil beam as a result of multiple scattering. The radius increases as the square of the depth, rather than the 3/2-power of depth dependence

characteristic of theories based on a small mean square scattering angle.

The second case concerns the degradation of the image of a point-source object emitting almost isotropically in angle. At the source, the intensity distribution is given by

$$\psi(\vec{r},\vec{\theta}) = \delta(\vec{r}) \frac{e^{-\theta^2/2(\Delta\theta)^2}}{2\pi(\Delta\theta)^2} \tag{21}$$

The angle $\Delta\theta$ has no intrinsic significance; it is chosen to satisfy $\theta_0 << \Delta\theta \lesssim 1$ so that the small-angle approximation remains valid while the initial distribution is diffuse compared to the typical scattering angle θ_0. Direct integration yields the initial Fourier transform

$$\Psi(\vec{k},\vec{p}) = e^{\frac{-p^2\theta_0^2}{2}} . \tag{22}$$

Starting with the initial condition (22), one can solve (9) by integrating along characteristics to obtain

$$\Psi(z,\vec{k},\vec{p}) = \exp\left\{-\tau_\alpha - \frac{s_0{}^2(\Delta\theta)^2}{2} + \int_0^z \beta_T(z')\left[X(s)-1\right]dz'\right\} \tag{23}$$

where

$$s_0 = |\vec{p} + \vec{k}z| \ . \tag{24}$$

Let us consider a case where the point-source is located at a given depth z, and where the image is viewed from overhead. Figure 2 makes it clear that this image depends on $\psi(z,\vec{r},\vec{\theta})$ evaluated at $\vec{\theta} = 0$. A Fourier inversion of (23) yields

$$\psi(z,\vec{r},\vec{\theta}) = \int \frac{d^2\vec{k}}{(2\pi)^2} e^{i\vec{k}\cdot\vec{r}} \int \frac{d^2\vec{p}}{(2\pi)^2} e^{i\vec{p}\cdot\vec{\theta}} \Psi(z,\vec{k},\vec{p}) \ , \tag{25}$$

the change of variables $\vec{p} = \vec{p}' - \vec{k}z$ leads to

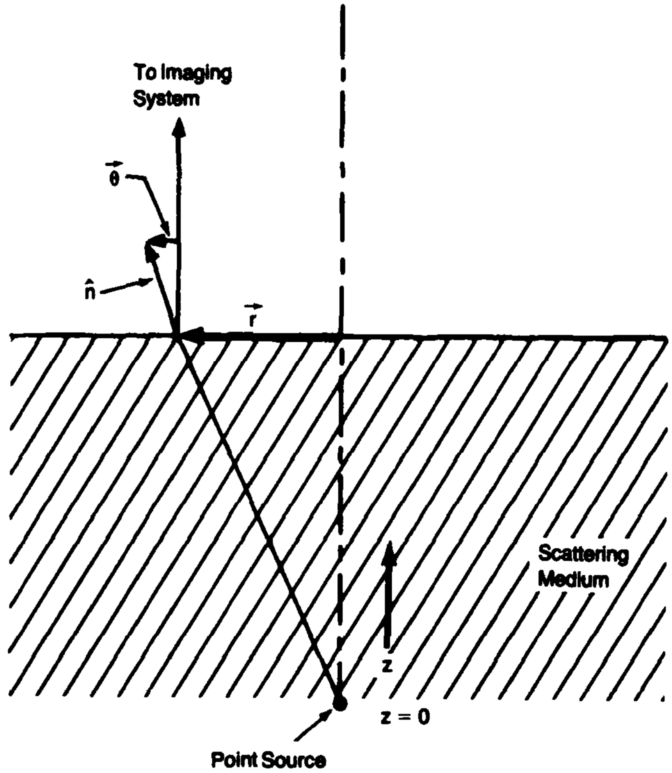
$$\psi(z,\vec{r},\vec{\theta}) = \int \frac{d^2\vec{k}}{(2\pi)^2} e^{i\vec{k}\cdot(\vec{r}-z\vec{\theta})} \int \frac{d^2\vec{p}'}{(2\pi)^2} e^{i\vec{p}'\cdot\vec{\theta}} \Psi(z,\vec{k},\vec{p}) \ . \tag{26}$$

Again, we will use the Taylor expansion of $X(s)$

$$\psi(z,\vec{r},\vec{\theta}) = e^{-\tau_\alpha} \int \frac{d^2\vec{k}}{(2\pi)^2} e^{i\vec{k}\cdot(\vec{r}-z\vec{\theta})} \int \frac{d^2\vec{p}'}{(2\pi)^2} \cdot \exp\left\{ \frac{-p'^2(\Delta\theta)^2}{2} \right.$$

$$\left. - \frac{2}{\pi}\,\theta_0 \int_0^z \beta_T(z')\ [p'^2 + k^2 z'^2 - 2\vec{p}'\cdot\vec{k}z']^{1/2} \right\} . \tag{27}$$

Figure 2. Geometry relevant to degradation of a point-source image. The origin of z is at the point source. The vector $\vec{\theta}$ is the component of a unit $\hat{n}$ vector which is perpendicular to the z-axis. The imaging system is far removed and records only radiation with $\vec{\theta} = 0$.



At this point, it is useful to regard $\Delta\Theta$ as formally large which permits one to ignore p' relative to kz' in the integral involving $\beta_T(z')$. Thus one obtains

$$\psi(z,\vec{r},\vec{\theta}) = \frac{e^{-\tau_\alpha}}{2\pi(\Delta\Theta)^2} \int \frac{d^2\vec{k}}{(2\pi)^2} \exp\left\{i\vec{k}\cdot\left(\vec{r}-z\vec{\theta}\right) - \frac{2\Theta_0 k}{\pi} \int_0^z \beta_T(z')z'dz'\right\}$$

$$= \frac{e^{-\tau_\alpha}}{2\pi(\Delta\Theta)^2} \frac{\bar{z}\Theta_0}{\pi^2\left[\left|\vec{r} - z\vec{\theta}\right|^2 + \left(\frac{2}{\pi}\bar{z}\;\Theta_0\right)^2\right]^{3/2}} \qquad (28)$$

where

$$\bar{z} = \int_0^z \beta_T(z')z'dz' \; . \qquad (29)$$

Note definitions (18) and (29) for $\bar{z}$ are identical when the different origin of z is taken into account. The factor $1/\left[2\pi(\Delta\Theta)^2\right]$ represents simply the intensity-per-steradian of the point source.

The image formed according to the system portrayed in Figure 2 depends on an evaluation of (28) at $\vec{\theta} = 0$. Hence, multiple scattering degrades the image of a point-source into one with the finite size $\Delta r \approx \left(2\bar{z}\Theta_0/\pi\right)$.

It is important to emphasize that our approximate methods are appropriate to the large optical depth regime

$$\tau = \int_0^z \beta_T(z')dz' > 1 . \tag{30}$$

But even in this regime, there will be an unscattered component of strength $e^{-\tau}$ which could appear much brighter in an imaging system if there is a true point source. In experiments dealing with the propagation of laser beams, this bright core dominates the experimental results[1,5], even though the fraction of photons in the core decreases exponentially. But for extended objects of finite surface brightness, the image will be dominated by the multiplely-scattered component which contains most of the intensity.

RANGE-STRAGGLING OF A LASER RADAR

The range resolution of a laser radar in a multiple-scattering medium is degraded because the photons do not travel in straight lines. Let us write the distribution function f as

$$f = f\left(z, z_2, \vec{r}, \vec{\theta}\right) \tag{31}$$

$$z_2 = \frac{ct}{n} - z \tag{32}$$

and include second-order multiple scattering effects in the radiative transfer equation

$$\frac{n}{c}\frac{\partial f}{\partial t} + \left(1 - \frac{\theta^2}{2}\right)\frac{\partial f}{\partial z} + \vec{\theta} \cdot \frac{\partial f}{\partial \vec{r}}$$
$$= -(\alpha + \beta_T) f + \beta_T \int d^2 \vec{\theta}' \; f(z, z_2, \vec{r}, \vec{\theta}') \; X(\vec{\theta} - \vec{\theta}') \; . \tag{33}$$

Range-straggling appears through the $\theta^2/2 \; \frac{\partial f}{\partial z}$ term. In any radar system,

$$\frac{\partial f}{\partial z_2} >> \frac{\partial f}{\partial z} \tag{34}$$

and multiple-scale arguments permit one to simplify (33) to

$$\frac{\partial f}{\partial z} + \frac{\theta^2}{2}\frac{\partial f}{\partial z_2} + \vec{\theta} \cdot \frac{\partial f}{\partial \vec{r}} = -(\alpha+\beta_T)f + \beta_T \int d^2\, \vec{\theta}' f(\vec{\theta}')\, X(\vec{\theta}-\vec{\theta}) \,. \tag{35}$$

Analytic progress can be made if we restrict our attention to $g(z,z_2,\vec{\theta})$ defined by

$$g(z,z_2,\vec{\theta}) = \int d^2\, \vec{r}\, f \,. \tag{36}$$

Physically, g_2 represents the distribution function governing a plane wave propagating vertically downward. Clearly g satisfies the equation

$$\frac{\partial g}{\partial z} + \frac{\theta^2}{2}\frac{\partial g}{\partial z_2} = -(\alpha+\beta_T)g + \beta_T \int d^2\, \vec{\theta}'\, g(\vec{\theta}')\, X(\vec{\theta}-\vec{\theta}') \tag{37}$$

which can be Fourier-transformed in angle to be

$$\frac{\partial G}{\partial z} - \frac{1}{2p}\frac{\partial}{\partial p}\, p\, \frac{\partial}{\partial p}\frac{\partial G}{\partial z_2} = \beta_T G(p)[X(p)-1] \tag{38}$$

where

$$G = e^{\tau_\alpha} \int d\vec{\theta}\, e^{-i\vec{p}\cdot\vec{\theta}}\, g(z,z_2,\vec{\theta}) \,. \tag{39}$$

The Taylor expansion of $X(p)$ results in

$$\frac{\partial G}{\partial z} - \frac{1}{2p}\frac{\partial}{\partial p}\, p\, \frac{\partial}{\partial p}\frac{\partial G}{\partial z_2} = -\beta_T \frac{2}{\pi}\, \Theta_0 pG \ . \tag{40}$$

Our delta-function initial conditions permit the use of similarity arguments when β_T is constant. Thus, we can introduce the similarity variables

$$u = p\beta_T 2\Theta_0 z/\pi$$

$$v = 2z_2\left[\left(\beta_T 2\Theta_0/\pi\right)^2 z^3\right]^{-1}$$

$$F = G\left[\left(\beta_T 2\Theta_0/\pi\right)^2 z^3\right] \tag{41}$$

and transform (40) into

$$3\frac{\partial}{\partial v}\left(vF\right) + \frac{1}{u}\frac{\partial}{\partial u}\, u\, \frac{\partial}{\partial u}\left(\frac{\partial F}{\partial v}\right) = uF + u\frac{\partial F}{\partial u} \ . \tag{42}$$

The function F will vanish for $v < 0$ and as $v \to \infty$. An integration with respect to v yields

$$3vF + \frac{1}{u}\frac{\partial}{\partial u}\, u\, \frac{\partial}{\partial u} F = u\left(1 + \frac{\partial}{\partial u}\right)H \tag{43}$$

$$H = \int_0^v F dv \ . \tag{44}$$

Along the line $v = 0$, the solution is $F = \text{constant}$. The constant will be determined by the normalization condition

$$\int_0^{\infty} F dv = e^{-u} \tag{45}$$

which is readily derivable from (42).

Further progress in solving (43 - 44) does not appear possible without computational work. But the similarity scaling (41) shows that the range resolution scales as

$$z_2 = \left(\frac{\beta_T 2\Theta_0}{\pi}\right)^2 z^3 . \tag{46}$$

An accurate computational solution of (42) depends on obtaining the correct starting values as $v \to 0$ and as $v \to \infty$. Near $v = 0$, we can use the trial solution

$$F = v^n F_1(y) \qquad y = u\, v^{1/3} . \tag{47}$$

Substitution of this form into equation (42) yields the result as $v \to 0$

$$\frac{1}{y}\frac{\partial}{\partial y}\, y \frac{\partial}{\partial y}\left(nF_1 + \frac{y}{3}\frac{\partial F_1}{\partial y}\right) = y\, F_1 \tag{48}$$

where the correction terms are of order $v^{1/3}$. For large y, the solution takes the form

$$F_1 = e^{-\frac{\eta}{2}}\left[A \sin\left(\frac{\sqrt{3}\eta}{2}\right) + B \cos\left(\frac{\sqrt{3}\eta}{2}\right)\right] \tag{49}$$

when

$$\eta = 3^{1/3} y \ .$$

In the limit $y \to 0$, the desired solution has the form

$$F_1 = a\left[1 + \frac{y^3}{9(n+1)} + \ . \ . \ . \ .\right]. \tag{50}$$

These three free parameters A, B, n permit one to solve for a, and to liminate the two divergent solutions near the origin.

In the limit $v \to \infty$, equation (42) is solved by

$$F = \frac{v^3 e^{-u}}{\left(vu^3\right)^s} \tag{51}$$

where we require $s > 1$ in order for the normalization condition (45) to converge. Let us note that (51) is not valid for $u \lesssim \frac{1}{v^{1/3}}$. We also require that $\int_0^\infty Fu \, du < \infty$, leading to $s < 5/3$.

SUMMARY

A simple fit to the shape of the differential cross section allows substantial analytic progress to be made, including equation (20) for the radial spreading of underwater intensity, equation (28) for the angle-radius relationship for an underwater point source of illumination, and estimate (46) for the range resolution of a laser radar.

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